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New Physics Effects in Doubly Cabibbo Suppressed D Decays

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Abstract

The most sensitive experimental searches for D^0–$\bar{D}^0$ mixing use $D^0 \to K^+\pi^-$ decays. It is often assumed that effects of New Physics and, in particular, CP violation, can appear through the mixing, while the $c \to du\bar{s}$ decay amplitude cannot have significant contributions from New Physics and is, therefore, CP conserving to a good approximation. We examine this assumption in two ways. First, we calculate the contributions to the decay in various relevant models of New Physics: Supersymmetry without R-parity, multi-scalar models, left-right symmetric models, and models with extra quarks. We find that phenomenological constraints imply that the New Physics contributions are indeed small compared to the standard model doubly Cabibbo suppressed amplitude. Second, we show that many of our constraints hold model-independently. We find, however, one case where the model-independent bound is rather weak and a CP violating contribution of order 30% is not excluded.

I. INTRODUCTION

The decay $\bar{D}^0 \to K^+\pi^-$ proceeds via the quark sub-process $\bar{c} \to d\bar{u}\bar{s}$ and is Cabibbo favored:

$$A^{\rm SM}(\bar{D}^0 \to K^+\pi^-) \propto G_F|V_{cs}V_{ud}|. \tag{1.1}$$

The decay $D^0 \to K^+\pi^-$ proceeds via the quark sub-process $c \to du\bar{s}$ and is doubly Cabibbo suppressed:

$$A^{\rm SM}(D^0 \to K^+\pi^-) \propto G_F|V_{cd}V_{us}|. \tag{1.2}$$

If D^0–$\bar{D}^0$ mixing is large, then there could be a significant contribution to the latter from $D^0 \to \bar{D}^0 \to K^+\pi^-$, where the second stage is Cabibbo favored. The most sensitive experimental searches for D^0–$\bar{D}^0$ mixing use indeed this process. The fact that the first-mix-then-decay amplitude gives a different time dependence than the direct decay allows experimenters to distinguish between the two contributions and to set unambiguous upper bounds on the mixing.

The standard model (SM) prediction for D^0–$\bar{D}^0$ mixing, $(\Delta m_D/m_D)_{\rm SM} \sim 10^{-16}$ [1–14], is well below the present experimental sensitivity, $(\Delta m_D/m_D)_{\rm exp} < 8.5 \times 10^{-14}$ [15–19]. If mixing is discovered within an order of magnitude of present bounds, its theoretical explanation will require contributions from New Physics. Even more convincing evidence for New Physics will arise if CP violation plays a role in the $D^0 \to K^+\pi^-$ decay [20,21]. The reason is that, while the calculation of the total rate suffers from large hadronic uncertainties related to the long distance contributions, the SM prediction that there is no CP violation is very safe since it is only related to the fact that the third generation plays almost no role in both the mixing and the decay.

Most if not all present analyses of the search for D^0–$\bar{D}^0$ mixing through $D \to K\pi$ decays make the assumption that the New Physics can affect significantly the mixing but not the decay. This is a plausible assumption. The SM contribution to the mixing is highly suppressed because it is second order in α_W and has a very strong GIM suppression factor, $m_s^4/(M_W m_c)^2$. The mixing is then sensitive to New Physics which could contribute at tree level (as in multi-scalar models), or through strong interactions (as in various supersymmetric models), etc. On the other hand, the SM contribution to the decay is through the tree-level W-mediated diagram. One does not expect that New Physics could give competing contributions.

Yet, since the decay in question is doubly Cabibbo suppressed, one may wonder if indeed the assumption that it gets no New Physics contributions is safe. It is the purpose of this work to test this assumption in a more concrete way. (For previous work on related processes, see [22–24].) We examine various reasonable extensions of the standard model with new tree level contributions to the decay. For each model, we present the relevant phenomenological constraints and find an upper bound on the new contributions to $D^0 \to K^+\pi^-$.

From (1.1) and (1.2) we get the following (naive) estimate for the ratio of amplitudes:

$$\left|\frac{A^{\rm SM}(D^0 \to K^+\pi^-)}{A^{\rm SM}(\bar{D}^0 \to K^+\pi^-)}\right| \sim \left|\frac{V_{cd}V_{us}}{V_{cs}V_{ud}}\right| \sim 0.05\,. \tag{1.3}$$

The value of this ratio from the recent CLEO results [19] is about 0.058. Thus, if New Physics contributions to $D^0 \to K^+\pi^-$ are to compete with the doubly Cabibbo suppressed SM amplitude, the corresponding effective New Physics coupling G_N should satisfy

$$G_N \gtrsim 10^{-2}\, G_F\,. \tag{1.4}$$

In Section II we investigate if this is possible in various New Physics scenarios. In Section III we study model-independent bounds on new tree-level contributions to $D^0 \to K^+\pi^-$. We conclude in Section IV.

II. SPECIFIC MODELS

A. Supersymmetry without R-parity

Supersymmetry without R-parity (R_p) predicts new tree diagrams contributing to the decay. The lepton number violating terms $\lambda'_{ijk}L_iQ_jd^c_k$ give a slepton-mediated contribution with an effective coupling:

$$G_N^{\lambda'} = \frac{\lambda'_{21k}\lambda'^*_{12k}}{4\sqrt{2}M^2(\tilde{\ell}^-_{Lk})}\,. \tag{2.1}$$

These couplings are severely constrained by K^0–$\bar{K}^0$ mixing (see *e.g.* [25]):

$$\lambda'_{21k}\lambda'^*_{12k} \lesssim 10^{-9} \quad (\text{for } M(\tilde{\ell}^-_{Lk}) = 100 \ \text{GeV}). \tag{2.2}$$

This rules out any significant contribution to $D^0 \to K^+\pi^-$ from slepton exchange in models of R_p violation:

$$\frac{G_N^{\lambda'}}{G_F|V_{cd}V_{us}|} \lesssim 3\times 10^{-8}\,. \tag{2.3}$$

The baryon number violating terms $\lambda''_{ijk}u^c_i d^c_j d^c_k$ give a squark-mediated contribution with an effective coupling

$$G_N^{\lambda''} = \frac{\lambda''_{113}\lambda''^*_{223}}{4\sqrt{2}M^2(\tilde{b}_R)}\,. \tag{2.4}$$

The λ''_{223} couplings is only constrained by requiring that it remains in the perturbative domain up to the unification scale and could be of order unity [26]. The λ''_{113} coupling is, however, severely constrained by the upper bound on $n-\bar{n}$ oscillations [27]:

$$|\lambda''_{113}| \lesssim 10^{-4} \quad (\text{for } M(\tilde{q}) = 100 \ \text{GeV}). \tag{2.5}$$

This rules out a significant contribution to $D^0 \to K^+\pi^-$ from squark exchange in models of R_p violation:

$$\frac{G_N^{\lambda''}}{G_F|V_{cd}V_{us}|} \lesssim 3\times 10^{-3}\,. \tag{2.6}$$

B. Multi-Scalar Models

Extensions of the scalar sector, beyond the single Higgs doublet of the SM, predict new tree diagrams contributing to the decay.

In two Higgs doublet models (2HDM) with natural flavor conservation, there is a charged Higgs ($H^\pm$) mediated contribution. The trilinear coupling of the physical charged Higgs to the $u_i\,\bar{d}_j$ bilinear is

$$- \mathcal{L}_{H^\pm} = \frac{ig}{\sqrt{2}m_W}\,\overline{u_i}\,\Big[m_{u_i}\cot\beta\,P_L + m_{d_j}\tan\beta\,P_R\Big]\,V_{ij}\,d_j H^+ + h.c.\,, \tag{2.7}$$

where m_W is the mass of the W-boson, m_q is the mass of the quark q, $\tan\beta = v_u/v_d$ is the ratio of vevs and $P_{R,L} = (1 \pm \gamma_5)/2$. It follows that the charged Higgs mediated contribution is also doubly Cabibbo suppressed. Then, for large $\tan\beta$, the suppression with respect to the SM contribution is given by

$$\frac{G_N^{H^\pm}}{G_F|V_{cd}V_{us}|} \simeq \frac{m_d\,m_s\tan^2\beta}{M^2_{H^\pm}} \lesssim 4\times 10^{-4}\,. \tag{2.8}$$

To obtain the upper bound, we used the constraint from $b \to c\,\tau\,\nu$ [28,29]:

$$\tan\beta \lesssim 0.5\left(\frac{M_{H^\pm}}{\mathrm{GeV}}\right)\,, \tag{2.9}$$

and the ranges of quark masses given in Ref. [30]. For $\tan\beta \simeq 1$ we have

$$\frac{G_N^{H^\pm}}{G_F|V_{cd}V_{us}|} \simeq \frac{m_s\,m_c}{M^2_{H^\pm}} \lesssim 10^{-4}\,. \tag{2.10}$$

To obtain the upper bound, we used $M_{H^\pm} \gtrsim 54.5$ GeV [30]. Thus there are no significant contributions to $D^0 \to K^+\pi^-$ from charged Higgs exchange within 2HDM.

Multi Higgs doublet models with natural flavor conservation but with more than two Higgs doublets have parameters that are less constrained and, in particular, provide new sources of CP violation. There are several charged scalars that can mediate the $D^0 \to K^+\pi^-$ decay. If we take the simplest case that only one of them contributes in a significant way (see *e.g.* [31]), then its couplings are similar to those of Eq. (2.7) except that $\tan\beta$ and $\cot\beta$ are replaced by, respectively, X and Y. In general, X and Y are complex and, moreover, $|XY| \neq 1$. Eq. (2.8) is modified:

$$\frac{G_N^{H^\pm}}{G_F|V_{cd}V_{us}|} \simeq \frac{m_d\,m_s|X|^2}{M^2_{H^\pm}} \lesssim 10^{-2}\,. \tag{2.11}$$

To obtain the upper bound, we used the perturbativity bound $|X| \lesssim 130$ [32,31] and the lower bound on $M_{H^\pm}$. Note that this contribution is not only constrained to be small, but also it carries no new CP violating phase. In contrast, the new contribution that replaces that of Eq. (2.10) does carry a new phase:

$$\frac{G_N^{H^\pm}}{G_F|V_{cd}V_{us}|} \simeq \frac{m_s\, m_c Y X^*}{M_{H^\pm}^2} \lesssim 3\times 10^{-4}\,. \tag{2.12}$$

To obtain the upper bound, we used the constraint from $b\to s\gamma$, $|XY|\lesssim 4$ [31]. The bound on a CP violating contribution is even somewhat stronger, since the measurement of $b\to s\gamma$ gives $\mathcal{I}m(X^*Y)\lesssim 2$ [33]. In any case, the contribution from charged Higgs exchange in multi Higgs doublet models is, at most, at the percent level. The CP violating part of this contribution is at most of order 10^{-4}.

It is possible that Yukawa couplings are naturally suppressed by flavor symmetries rather than by natural flavor conservation [6]. In such a framework, there is a contribution to $D^0\to K^+\pi^-$ from neutral scalar exchange. To estimate these contributions, we use the explicit models of Ref. [34]. Here, a horizontal $U(1)_{\mathcal{H}}$ symmetry is imposed. At low energies, the symmetry is broken by a small parameter λ (usually taken to be of the order of the Cabibbo angle, $\lambda\sim 0.2$), leading to selection rules. The scalar sector consists of two Higgs doublets, ϕ_u and ϕ_d, and a single scalar singlet S. The effective coupling of the S scalar to quarks is given by

$$-\mathcal{L}_S = Z_{ij}^q\, S\, \overline{q_{iR}}\, q_{jL} + h.c. \quad (q=u,d,\quad i,j=1,2,3). \tag{2.13}$$

The order of magnitude of Z_{ij}^q is determined by the selection rules related to the broken flavor symmetry:

$$Z_{ij}^q \sim \frac{M_{ij}^q}{\langle S\rangle}\,,\qquad M_{ij}^q \sim \lambda^{\mathcal{H}(q_{jL})+\mathcal{H}(\overline{q_{iR}})+\mathcal{H}(\phi_q)}\langle\phi_q\rangle\,. \tag{2.14}$$

The horizontal charges $\mathcal{H}$ of the quark and Higgs fields are determined by the physical flavor parameters:

$$\begin{aligned}|V_{ij}| &\sim \lambda^{\mathcal{H}(q_{iL})-\mathcal{H}(q_{jL})},\\ m(q_i) &\sim \lambda^{\mathcal{H}(q_{iL})+\mathcal{H}(\overline{q_{iR}})+\mathcal{H}(\phi_q)}\langle\phi_q\rangle.\end{aligned} \tag{2.15}$$

Using (2.15) we can express the suppression of the relevant Yukawa couplings in terms of the quark masses and mixing angles:

$$|Z_{uc}^u| \sim \frac{m_c|V_{12}|}{\langle S\rangle}\,,\quad |Z_{cu}^u| \sim \frac{m_u}{\langle S\rangle|V_{12}|}\,,\quad |Z_{ds}^d| \sim \frac{m_s|V_{12}|\tan\beta}{\langle S\rangle}\,,\quad |Z_{sd}^d| \sim \frac{m_d\tan\beta}{\langle S\rangle|V_{12}|}\,. \tag{2.16}$$

These couplings give rise to various operators that induce $c\to u d\overline{s}$ at tree level. For the leading contributions, we find

$$\frac{G_N^S}{G_F|V_{cd}V_{us}|} \sim \frac{m_c m_s\tan\beta}{\langle S\rangle^2}\frac{m_W^2}{m_S^2} \lesssim 5\times 10^{-3}\,. \tag{2.17}$$

To obtain the upper bound, we used $\tan\beta\lesssim 130$ and the very conservative bound $\langle S\rangle\sim m_S > m_W$. Other models [6,35,36] give a similar or even stronger suppression. We conclude that there are no significant contributions to $D^0\to K^+\pi^-$ from neutral Higgs exchange within multi-scalar models with approximate flavor symmetries.

C. Left-Right Symmetric Models

Left-right symmetric (LRS) models predict new tree-level contributions, mediated by the W_R gauge bosons. The relevant interactions are given by

$$-\mathcal{L}_{CC} = \frac{g_R}{\sqrt{2}}\,\overline{u_{iR}}\,\gamma_\mu V^R_{ij}\,d_{jR}W_R^{\mu+} + h.c.\,, \tag{2.18}$$

where V^R is the mixing matrix for the right-handed quarks. For a general model of an extended electroweak gauge group $G = SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, the interactions of Eq. (2.18) lead to

$$\frac{G_N^{W_R}}{G_F\,|V_{cd}V_{us}^*|} = \frac{g_R^2}{g_L^2}\frac{m_{W_L}^2}{m_{W_R}^2}\left|\frac{V_{cd}^R V_{us}^{R*}}{V_{cd}V_{us}^*}\right|\,. \tag{2.19}$$

However, in left-right symmetric models, an extra discrete symmetry is imposed. It leads to the relation $g_L = g_R$ and, in models of spontaneous CP violation or of manifest left-right symmetry, to $|V_{ij}| = |V_{ij}^R|$. Then Eq. (2.19) is simplified:

$$\frac{G_N^{\rm LRS}}{G_F\,|V_{cd}V_{us}^*|} = \frac{m_{W_L}^2}{m_{W_R}^2} \lesssim \frac{1}{430}\,, \tag{2.20}$$

where the upper bound comes from the Δm_K constraint [37].

In $SU(2)_L\times SU(2)_R\times U(1)_{B-L}$ models where V and V_R are independent mixing matrices, it is possible to avoid the ΔM_K constraints [38,39]. This is done by fine tuning the relevant entries in V_R to be very small. In particular, it was shown that in such a framework there could be interesting implications on CP violation in the B system [39]. However, as concerns the $D^0 \to K^+\pi^-$ decay, the situation is different: the *same* mixing elements that contribute to $D^0 \to K^+\pi^-$, that is $V_{cd}^R V_{us}^{R*}$, contribute also to $K-\bar{K}$ mixing. If they are switched off, to avoid the Δm_K constraint, the new contribution to $D^0 \to K^+\pi^-$ vanishes as well. One can see that independently of the details of the model by noticing that the $G_N^{W_R}$ effective coupling of Eq. (2.19) can be combined with the flavor-changing $G_F V_{cd}V_{us}^*$ coupling of the SM to produce a contribution to $K-\bar{K}$ mixing. Indeed, one finds for the CP conserving contribution [38]:

$$\mathcal{R}e\left(\frac{G_N^{W_R}}{G_F V_{cd}V_{us}^*}\right) \lesssim 0.2\,, \tag{2.21}$$

and for the CP violating contribution [39]:

$$\mathcal{I}m\left(\frac{G_N^{W_R}}{G_F V_{cd}V_{us}^*}\right) \lesssim 0.002\,. \tag{2.22}$$

We learn that in such fine-tuned models, the W_R-mediated contribution to the decay rate could be non-negligible, but the CP violating contribution is very small.

D. Extra Quarks in SM Vector-Like Representations

In models with non-sequential ('exotic') quarks, the Z-boson has flavor changing couplings, leading to a Z-mediated contribution to the $D^0 \to K^+\pi^-$ decay. For example, in models with additional up quarks in the vector-like representation $(\mathbf{3},\mathbf{1},+2/3)\oplus(\bar{\mathbf{3}},\mathbf{1},-2/3)$ and additional down quarks in the vector-like representation $(\mathbf{3},\mathbf{1},-1/3)\oplus(\bar{\mathbf{3}},\mathbf{1},+1/3)$, the flavor changing Z couplings have the form

$$-\mathcal{L}_Z = \frac{g}{2\cos\theta_W}\left(U^u_{ij}\,\overline{u}_{Li}\gamma_\mu u_{Lj} - U^d_{ij}\,\overline{d}_{Li}\gamma_\mu d_{Lj}\right) Z^\mu + h.c.\,. \tag{2.23}$$

Here, $U^q = V_L^{q\dagger}\,\mathrm{diag}(1,1,1,0)\,V_L^q$, where V_L^q is the 4×4 diagonalizing matrix for $M_q M_q^\dagger$ (M_q being the quark mass matrix). The flavor changing couplings are constrained by ΔM_K and ΔM_D:

$$|U^d_{sd}| \lesssim 2\times 10^{-4}\,, \quad |U^u_{cu}| \lesssim 7\times 10^{-4}\,. \tag{2.24}$$

The resulting effective four fermi coupling is given by

$$\frac{G_N^Z}{G_F|V_{cd}V_{us}|} \simeq \frac{|U^d_{sd}U^u_{cu}|}{|V_{cd}V_{us}|} \lesssim 3\times 10^{-6}\,. \tag{2.25}$$

The same bound applies for the case of vector-like quark doublets, $(\mathbf{3},\mathbf{2},+1/6)\oplus(\bar{\mathbf{3}},\mathbf{2},-1/6)$. The flavor changing Z couplings are to right-handed quarks, with a mixing matrix $U^q = V_R^{q\dagger}\,\mathrm{diag}(0,0,0,1)\,V_R^q$. Here V_R^q is the 4×4 diagonalizing matrix for $M_q^\dagger M_q$.

We learn that a significant contribution to $D^0 \to K^+\pi^-$ from Z-mediated flavor changing interactions is ruled out.

III. MODEL INDEPENDENT ANALYSIS

We have seen that the contributions to $D^0 \to K^+\pi^-$ in various reasonable extensions of the SM cannot compete with the W-mediated process. Still, it would be useful if one could show *model-independently* that CP violation in decay can be neglected. We try to accomplish this task for all possible tree level contributions to the $D^0 \to K^+\pi^-$ decay. Our analysis proceeds as follows [40]: We first list all relevant (anti)quark bilinears and their transformation properties under the SM gauge group $\mathcal{G}_{SM} = SU(3)_C \times SU(2)_L \times U(1)_Y$. If the two quarks have the same (opposite) chirality, they couple to a scalar (vector) boson. Altogether there are ten possible bilinears (plus their hermitian conjugates) that are shown in Tab. 1 . Here Q denotes the left-handed quark doublet, $q = u, d$ denote the right-handed quark singlets, and the superscript c refers to the respective antiquarks. The examples given in the last column refer to the models discussed in Section II.

In general, the presence of a heavy boson $\mathcal{B}$ that couples to any of the above quark bilinears B_{ij} with trilinear couplings $\lambda^{\mathcal{B}}_{ij}$, where $i, j = 1, 2, 3$ refer to the quark flavors, gives rise to the four quark operator $B^\dagger_{ij}B_{kl}$ with the effective coupling

$$G_N^{\mathcal{B}} = C_{CG}\frac{\lambda_{ij}^{\mathcal{B}*}\lambda_{kl}^{\mathcal{B}}}{4\sqrt{2}M_{\mathcal{B}}^2}\,, \tag{3.1}$$

at energy scales well below the mass $M_{\mathcal{B}}$. (C_{CG} is the appropriate Clebsch-Gordan coefficient.) For intermediate diquarks, we only discuss color triplets. The discussion of color sextets follows similar lines.

Bilinear B	$SU(3)_C$	$SU(2)_L$	Y	Couples to Boson $\mathcal{B}$	Example
Qd^c	$\mathbf{1}$	$\mathbf{2}$	$1/2$	$\mathcal{S}(\mathbf{1},\mathbf{2},-1/2)$	$\tilde{L}$ (SUSY $\not{R}_p$)
$u^c d^c$	$\bar{\mathbf{3}}\otimes\bar{\mathbf{3}}$	$\mathbf{1}$	$-1/3$	$\mathcal{S}(\bar{\mathbf{3}},\mathbf{1},1/3)$	$\tilde{d}^c$ (SUSY $\not{R}_p$)
Qu^c	$\mathbf{1}$	$\mathbf{2}$	$-1/2$	$\mathcal{S}(\mathbf{1},\mathbf{2},1/2)$	H_u (2HDM)
QQ	$\mathbf{3}\otimes\mathbf{3}$	$\mathbf{2}\otimes\mathbf{2}$	$1/3$	$\mathcal{S}(\mathbf{3},\boldsymbol{L},-1/3)$ [$\boldsymbol{L}$=$\mathbf{1}$,$\mathbf{3}$]	
ud^c	$\mathbf{1}$	$\mathbf{1}$	1	$\mathcal{V}(\mathbf{1},\mathbf{1},-1)$	W_R (LRS)
qq^c	$\mathbf{1}$	$\mathbf{1}$	0	$\mathcal{V}(\mathbf{1},\mathbf{1},0)$	
Qd	$\mathbf{3}\otimes\mathbf{3}$	$\mathbf{2}$	$-1/6$	$\mathcal{V}(\mathbf{3},\mathbf{2},1/6)$	
Qu	$\mathbf{3}\otimes\mathbf{3}$	$\mathbf{2}$	$5/6$	$\mathcal{V}(\mathbf{3},\mathbf{2},-5/6)$	
QQ^c	$\mathbf{1}$	$\mathbf{2}\otimes\mathbf{2}$	0	$\mathcal{V}(\mathbf{1},\boldsymbol{L},0)$ [$\boldsymbol{L}$=$\mathbf{1}$,$\mathbf{3}$]	Z (extra q's)

Tab. 1: Quark-(Anti)Quark Bilinears

In order to predict the rates of the relevant hadronic process one would need to take into account QCD corrections as well as the hadronic matrix elements. Since we are mainly interested in ratios between the rates due to New Physics and those from the SM, using (3.1) is sufficient to obtain an order-of-magnitude estimate for such ratios.

The first entry in Tab. 1 is realized in supersymmetric models without R_p (SUSY $\not{R}_p$): $\mathcal{S}(\mathbf{1},\mathbf{2},-1/2)$ is the slepton doublet $\tilde{L}_k$, with $\lambda_{ij}^{\mathcal{S}(\mathbf{1},\mathbf{2},-1/2)} = \lambda'_{ijk}$. As we have pointed out in Section II A, non-vanishing $\lambda_{12}^{\mathcal{S}(\mathbf{1},\mathbf{2},-1/2)}$ and $\lambda_{21}^{\mathcal{S}(\mathbf{1},\mathbf{2},-1/2)}$ give rise not only to tree-level contributions to $D^0 \to K^+\pi^-$ but also to K^0–$\bar{K}^0$ mixing, which severely constraints the effective coupling G_N. In this case the bound arises only from the presence of the trilinear couplings and supersymmetry does not play a role. The bound in (2.3) is then model-independent.

The second entry in Tab. 1 is also realized in supersymmetric models without R_p: $\mathcal{S}(\mathbf{3},\mathbf{1},1/3)$ is the down squark $\tilde{d}^c_k$, with $\lambda_{ij}^{\mathcal{S}(\mathbf{3},\mathbf{1},1/3)} = \lambda''_{ijk}$. For the λ'' coupling, however, the constraint comes from the upper bound on $n-\bar{n}$ oscillations: to violate baryon number but conserve strangeness or beauty, an internal loop with charginos is required [27]. Supersymmetry does play a role in the bound on λ''_{113}, and the bound does not hold for a generic $\lambda_{11}^{\mathcal{S}(\mathbf{3},\mathbf{1},1/3)}$. More generally, there is no strong model-independent bound on any diagonal $\lambda_{ii}^{\mathcal{S}(\mathbf{3},\mathbf{1},1/3)}$ coupling. The bound on the scale of compositeness [30], $\Lambda(qqqq) \gtrsim 1.6$ TeV, suggests a bound for the $i=1$ case, $|\lambda_{11}^{\mathcal{S}(\mathbf{3},\mathbf{1},1/3)}| \lesssim 0.2$ which implies $G_N^{\mathcal{S}(\mathbf{3},\mathbf{1},1/3)} \lesssim 0.3\ G_F$ (assuming $|\lambda_{22}^{\mathcal{S}(\mathbf{3},\mathbf{1},1/3)}| \sim 1$). We learn then that one could construct models which incorporate color-triplet weak-singlet scalars where there is a large CP violating contribution to $D^0 \to K^+\pi^-$.

The coupling of Qu^c to $\mathcal{S}(\mathbf{1},\mathbf{2},1/2)$ appears in the two Higgs doublet model with natural flavor conservation, as discussed in Section II B. In this model, the effective coupling is suppressed by the quark masses and the CKM matrix elements. But also if the doublet $\mathcal{S}(\mathbf{1},\mathbf{2},1/2)$ is unrelated to the generation of the quarks masses, one can derive a model-independent bound, which only relies on the $SU(2)_L$ symmetry: Non-vanishing $\lambda_{12}^{\mathcal{S}(\mathbf{1},\mathbf{2},1/2)}$ and $\lambda_{21}^{\mathcal{S}(\mathbf{1},\mathbf{2},1/2)}$ give not only a charged scalar mediated contribution to $D^0 \to K^+\pi^-$, but also a neutral scalar mediated contribution to D^0–$\bar{D}^0$ mixing. We are assuming that the New Physics takes place at a scale that is comparable to or higher than the electroweak breaking scale, so that $SU(2)_L$ breaking effects are not large and the masses of the charged and neutral scalars are similar [41]. Consequently, the upper bound on D^0–$\bar{D}^0$ mixing translates into

$$G_N^{\mathcal{S}(\mathbf{1},\mathbf{2},1/2)_{12}} = \frac{\lambda_{12}^{\mathcal{S}(\mathbf{1},\mathbf{2},1/2)^*}\lambda_{21}^{\mathcal{S}(\mathbf{1},\mathbf{2},1/2)}}{4\sqrt{2}M^2_{\mathcal{S}(\mathbf{1},\mathbf{2},1/2)}} \lesssim 10^{-7}\ G_F\,, \tag{3.2}$$

too small to compete with the SM contribution.

The coupling of the QQ bilinear to a scalar field could induce $D^0 \to K^+\pi^-$ if the flavor diagonal entries, $\lambda_{11}^{\mathcal{S}(\mathbf{3},\boldsymbol{L},-1/3)}$ and $\lambda_{22}^{\mathcal{S}(\mathbf{3},\boldsymbol{L},-1/3)}$, are non-zero. For an $SU(2)_L$ singlet ($\boldsymbol{L}=\mathbf{1}$), the coupling is flavor anti-symmetric and therefore $\lambda_{ii}^{\mathcal{S}(\mathbf{3},\mathbf{1},-1/3)}=0$. For an $SU(2)_L$ triplet ($\boldsymbol{L}=\mathbf{3}$), the coupling is flavor symmetric and $\lambda_{ii}^{\mathcal{S}(\mathbf{3},\mathbf{3},-1/3)} \neq 0$ is possible. (For scalar $SU(3)_C$ sextets the situation would be reversed.) However, while the $Q_{EM}=-1/3$ component mediates $D^0 \to K^+\pi^-$, the $Q_{EM}=+2/3$ component induces K^0–$\bar{K}^0$ mixing and the $Q_{EM}=-4/3$ component induces D^0–$\bar{D}^0$ mixing. We find:

$$G_N^{\mathcal{S}(\mathbf{3},\mathbf{3},-1/3)_{12}} = \frac{\lambda_{11}^{\mathcal{S}(\mathbf{3},\mathbf{3},-1/3)}\lambda_{22}^{\mathcal{S}(\mathbf{3},\mathbf{3},-1/3)}}{4\sqrt{2}M^2_{\mathcal{S}(\mathbf{3},\mathbf{3},-1/3)}} \lesssim 10^{-7}\ G_F\,, \tag{3.3}$$

too small to compete with the SM contribution.

Among the vector bosons listed in Tab. 1 we already encountered specific examples for the color singlets $\mathcal{V}(\mathbf{1},\mathbf{1},-1)$ (W_R^- in LRS models) and $\mathcal{V}(\mathbf{1},\mathbf{3},0)$ (Z-induced FCNCs due to extra quarks). The discussion we presented in Section II C can be generalized to any theory that contains a vector boson $\mathcal{V}(\mathbf{1},\mathbf{1},-1)$ that couples to ud^c as in (2.18). Note that the W_R (being a gauge boson) has flavor diagonal couplings in the flavor basis and only the charged components induce flavor transitions between the mass eigenstates, while the neutral component cannot mediate FCNCs. Still, as we have seen in Section II C, the contribution from the left-right box diagram to ΔM_K and ϵ_K imposes severe constraints on the $D^0 \to K^+\pi^-$ amplitude due to $\mathcal{V}(\mathbf{1},\mathbf{1},-1)$ exchange and rules out significant CP violation in the decay.

For the generic coupling $\lambda^{\mathcal{V}(\mathbf{1},\mathbf{3},0)}$ we can adopt the specific result obtained in Section II D. Since the argument based on the bounds from K^0–$\bar{K}^0$ and D^0–$\bar{D}^0$ oscillations only used the trilinear couplings (2.23), it can be generalized to the generic couplings $\lambda^{\mathcal{V}(\mathbf{1},\mathbf{1},0)}$. Because all quark-antiquark bilinears B_{ij} that couple to $\mathcal{V}(\mathbf{1},\mathbf{1},0)$ are gauge-invariant one can easily

see that $\mathcal{V}(\mathbf{1},\mathbf{1},0)$ exchange induces not only the flavor-conserving effective operator $B_{ij}^\dagger B_{ij}$ but also the flavor-violating operator $B_{ij}B_{ij}$ that gives rise to neutral meson mixing.

For the remaining vector couplings in Tab. 1, the decay $D^0 \to K^+\pi^-$ can be induced, if the flavor diagonal couplings $\lambda_{ii}^{\mathcal{V}(\mathbf{3},\mathbf{2},Y)}$ for both $i = 1$ and 2 are non-zero. Note that the intermediate vector boson carries color. Since the respective quark bilinears contain one $SU(2)_L$ doublet the effective operator that gives rise to $D^0 \to K^+\pi^-$ is related by an $SU(2)_L$ rotation to an operator that induces K^0–$\bar{K}^0$ [for $\mathcal{V}(\mathbf{3},\mathbf{2},1/6)$] and D^0–$\bar{D}^0$ [for $\mathcal{V}(\mathbf{3},\mathbf{2},-5/6)$] oscillations at tree level. Since $SU(2)_L$ breaking effects are small [41], the data from neutral meson mixing imply $G_N^{\mathcal{V}(\mathbf{3},\mathbf{2},1/6)} \lesssim 10^{-7}\ G_F$, and $G_N^{\mathcal{V}(\mathbf{3},\mathbf{2},-5/6)} \lesssim 10^{-6}\ G_F$, ruling out any significant contribution to $D^0 \to K^+\pi^-$.

IV. CONCLUSIONS

We have examined well-motivated extensions of the standard model that give new, tree-level contributions to the $D^0 \to K^+\pi^-$ decay. We showed that in all the models that we considered, strong phenomenological constraints imply that these contributions can be safely neglected.

We have extended our discussion to a model-independent analysis of all possible tree level contributions to the decay. We found that there is only one case where a large contribution to $D^0 \to K^+\pi^-$ is possible. This is the case where two right-handed quarks, $u^c d^c$, couple to an $SU(2)_L$-singlet scalar, $\mathcal{S}(\bar{\mathbf{3}},\mathbf{1},1/3)$. Such a coupling is present in SUSY without R_p but in this model the relevant coupling is constrained by $n-\bar{n}$ oscillations, ruling out a contribution that is comparable to the SM doubly Cabibbo suppressed diagram.

In our analysis, we have implicitly assumed that there are no significant accidental cancellations between various contributions to the processes from which we derive our constraints. It is possible to construct fine-tuned models where there is a large new contribution to $D^0 \to K^+\pi^-$ while the related contributions to flavor changing neutral current processes are small.

We conclude that, in general, the assumption that New Physics effects could affect the $D^0 \to K^+\pi^-$ decay and, in particular, its CP violating part, only through D^0–$\bar{D}^0$ mixing is a good assumption and it holds to better than one percent in all the reasonable and well-motivated extensions of the standard model that we have examined. One could construct, however, viable (even if unmotivated) models where there is a new, large [$\mathcal{O}(0.3)$] and possibly CP violating contribution to the decay.

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